

# Dynamics of Chiral Symmetry Breaking in Nuclear Collisions

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## ABSTRACT

Measurements of disoriented chiral condensates in heavy ion collisions at RHIC and the LHC can yield fundamental information on the nature of the QCD phase transition. I review theoretical efforts to understand the evolution of the condensate and present new results on experimental signals in the single pion spectrum and in pion interferometry.

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The order of the chiral restoration phase transition in high temperature QCD is currently unknown. If the transition is second order, Rajagopal and Wilczek (RW) propose that the nonequilibrium dynamics in ion–ion collisions at RHIC can generate transient domains in which a macroscopic pion field develops.<sup>1</sup> These domains are called disoriented chiral condensates (DCC). Bjorken, Kowalski, Taylor and others pointed out that DCCs can lead to fluctuations in the charged and neutral pion spectra.<sup>2</sup> Hints of this behavior may have been seen in the Centauro cosmic ray events.<sup>3</sup> In ion–ion collisions, the ability of experimenters to identify DCCs amidst the background produced by conventional particle production and scattering critically depends on the domains’ size and energy content.

In this talk, I explore the RW mechanism for domain formation to determine the size of domains in nuclear collisions, reporting on work with Gocksch and Pisarski<sup>4</sup> and Müller<sup>5</sup>. I then briefly speculate on phenomenological consequences of DCC formation. [For a discussion of  $J/\psi$  suppression as presented at this meeting, see Ref. [6].]

Equilibrium high temperature QCD manifests a chiral symmetry if the light up and down quarks are taken to be massless. However, a phase transition occurs at a critical temperature  $T_c \sim 140$  MeV in which chiral symmetry is broken by the formation of a scalar  $\langle \bar{q}q \rangle$  condensate. In the real world chiral symmetry is explicitly broken at all temperatures by the few–MeV current quark masses. In that case the nature of the transition has yet to be established.<sup>7</sup> [For reviews of lattice and finite temperature QCD, see e.g. Refs. [8].]

Rajagopal and Wilczek pointed out that the chiral condensate can become temporarily disoriented in the nonequilibrium conditions encountered in heavy ion collisions. Near  $T_c$ , the approximate chiral symmetry implies that the scalar condensate is nearly equivalent to a pion–like pseudoscalar isovector condensate  $\sim \langle \bar{q}\gamma_5 \vec{\tau} q \rangle$ , where  $\vec{\tau}$  are the Pauli isospin matrices. Consequently, domains containing a macroscopic pion field can appear as the temperature drops below  $T_c$ . Such domains will eventually disappear as the system evolves towards the true vacuum in which only the scalar condensate is nonzero. In the finite–sized heavy ion system, however, the evolving domains can radiate pions preferentially according to their isospin content.

The excitation of a pion field over a substantial fraction of the collision volume can result in novel fluctuations in the number of neutral and charged pions. Consider a single ideal domain in which the pion field is oscillating along a fixed isospin direction  $\vec{\pi}$ . Pion production from this domain is proportional to the square of the field strength, so that the fraction of neutral pions  $f$  is

$$f = |\pi^0|^2 / \sum_{0,+,-} |\pi^a|^2 \equiv \cos^2 \theta. \quad (1)$$

All orientations of  $\vec{\pi}$  are equally likely, so that the probability of a finding a fraction  $f$  is

$$P(f) \propto d(\cos \theta) / df \propto f^{-1/2} \quad (2)$$

It is therefore most probable for a single domain to emit more charged pions than neutral.<sup>2</sup>

The detectability of domains in a nuclear collision depends on their size.<sup>4</sup> If the interaction volume in a nuclear collision is dominated by a single domain, then DCC formation can lead to a measurable isospin asymmetry. If a nuclear collision instead produces many

uncorrelated small domains, the spectrum of fluctuations would be gaussian, as expected, *e.g.*, if the pions were produced independently.

To develop a more concrete picture of how the chiral condensate evolves, we use the linear sigma model, in which the pion field is coupled to a scalar  $\sigma$  field that characterizes the scalar condensate<sup>1</sup>. The fields interact through a potential

$$V = \lambda (\vec{\pi}^2 + \sigma^2 - v^2)^2 / 4 - H\sigma. \quad (3)$$

Chiral symmetry transformations correspond to  $O(4)$  rotations of the vector field  $\Phi = (\sigma, \vec{\pi})$ . For  $H = 0$  this symmetry is spontaneously broken at zero temperature, with  $\langle \sigma \rangle = v$ . Current algebra implies that  $\langle \sigma \rangle = f_\pi = 93$  MeV.<sup>9</sup> The external field  $H$  ( $\propto$  the light quark masses) breaks the  $O(4)$  symmetry explicitly and gives the pions a small mass  $m_\pi = \sqrt{H/f_\pi} = 140$  MeV. In comparison, the sigma mass  $m_\sigma \sim \sqrt{2\lambda v^2}$  is quite large, perhaps  $\sim 600$  MeV.

Following Ref. [1], I will take this model with the parameters fixed at  $T = 0$  as the basis for a Ginzburg–Landau description of the chiral transition. For  $H = 0$ , the transition would then be second order with an  $O(4)$  order parameter  $\Phi$ . DCC domains would then be well defined entities in the thermodynamic limit. While there is strictly no transition for  $H \neq 0$ , *transient* domains can form as follows.

In the idealized ‘quench’ scenario proposed by Rajagopal and Wilczek, one assumes that the system initially has  $\Phi = 0$ , as appropriate at high temperature, and follows the development of the system using  $T = 0$  equations of motion derived from  $V$ . The initial state of the system is clearly unstable. The system “rolls down” from the unstable local maximum of  $V(\Phi)$  towards the nearly stable values with  $|\Phi| = v$  (the symmetry breaking term  $-H\sigma$  is relatively small). Field configurations with  $\vec{\pi} \neq 0$  develop during the roll-down period. The field will eventually settle into stable oscillations about the unique vacuum  $(f_\pi, \vec{0})$  for  $H \neq 0$ , but oscillations continue until interactions eventually damp the motion.

To be more concrete, the linearized equations of motion for the Fourier components of the pion field are:

$$\frac{d^2}{dt^2} \vec{\pi}_{\vec{k}} = \{\lambda v^2 - k^2\} \vec{\pi}_{\vec{k}}. \quad (4)$$

Field configurations with  $\langle \Phi \rangle = 0$  and momentum  $k < \sqrt{\lambda}v$  are unstable and grow exponentially; modes with higher momenta do not grow. The  $k = 0$  mode grows the fastest, with a time scale

$$\tau_R = \{\lambda v^2\}^{-1/2} \sim \sqrt{2}/m_\sigma. \quad (5)$$

The unstable modes grow for a time of order  $\tau_R$  until  $\Phi^2$  approaches  $v^2$ , after which the system oscillates about  $\sigma = f_\pi$ .

Significantly, the time during which growth can occur is quite short,  $\tau_R \sim 0.5$  fm for  $m_\sigma \sim 600$  MeV. Rajagopal and Wilczek found that the power  $\propto \langle \pi_{\vec{k}}^a \pi_{-\vec{k}}^a \rangle$  in the low momentum pion modes grows when the exact classical equations of motion are integrated, demonstrating that domains indeed form. [Here  $\langle \dots \rangle$  represents an average over the initial of the fields, which are taken to have a thermal spectrum.] What is not clear, however, is

whether the time during which the system is unstable is sufficient for truly large domains to form.

Gocksch, Pisarski and I<sup>4</sup> studied the domain size by numerically integrating the equations of motion in the quench scenario and extracting the spatial correlation function:

$$\langle \pi(\vec{x}, t) \pi(0, t) \rangle \propto \int d^3k \langle \pi_{-\vec{k}} \pi_{\vec{k}} \rangle e^{-i\vec{k} \cdot \vec{x}}, \quad (6)$$

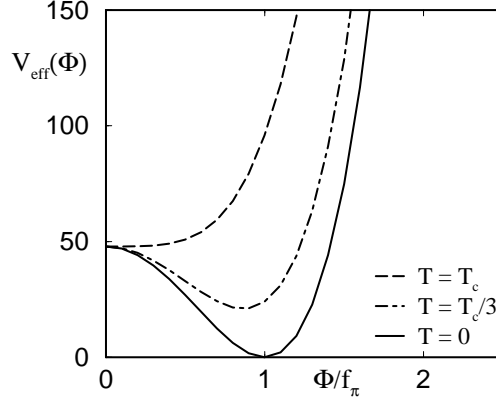
where isospin labels are implicit. This correlation function contains size information since  $\langle \pi(\vec{x}, t) \pi(0, t) \rangle \approx \langle \pi(0, t) \rangle^2$  holds inside the domain. The numerical results in Ref. [4] can be understood by taking  $\pi_k(t)$  from the linearized equation (4) and assuming that the fluctuations are initially thermal, so that  $\langle \pi_{-\vec{k}} \pi_{\vec{k}} \rangle \sim \{E(\exp E/T_c - 1)\}^{-1}$  for  $E = \{k^2 + m_\pi^2\}^{1/2}$ . A saddle-point integration yields  $\langle \pi(\vec{x}) \pi(0) \rangle \propto e^{-x^2/8\tau_R t}$ . The domain size,  $R_D \sim 2\tau_R \propto m_\sigma^{-1}$ , is set by the sigma mass. For  $m_\sigma = 600$  MeV, we find that domains are essentially pion sized. Numerical simulations confirm this result.

The lesson drawn from the quench scenario is that large nucleus-sized domains can occur only if the sigma mass is small.<sup>4,10,11</sup> However, it is quite possible that the effective  $m_\sigma$  is reduced in the high energy density heavy ion environment. In fact, at a truly second order phase transition we expect  $m_\sigma$  — the inverse correlation length — to vanish! In that case, the parameters of the effective potential for our Ginzburg-Landau model cannot be *naively* taken from zero-temperature physics as we have done.

To explore the role of the medium in domain formation, Müller and I studied the evolution of the condensate in the presence of a nonequilibrium bath of quasiparticles.<sup>5</sup> We find that larger domains are possible due to two effects, *i)* the reduction of  $m_\sigma$  and *ii)* the “annealing” of the system, rather than quenching, by the slowly expanding heavy ion system. [Our use of the word “annealing” is meant to distinguish our scenario from RW’s “quench” — the similarity of these terms to those used in condensed matter physics is in some respects misleading.]

I illustrate the role of mass reduction by considering a system near equilibrium. The system is described by an effective potential  $V_{\text{eff}}$  that has the behavior shown in Fig. 1.<sup>12</sup> In the Hartree mean-field approximation, the change of  $V_{\text{eff}}(\Phi)$  as a function of temperature is determined essentially by the coefficient of  $\Phi^2$ , *i.e.* the mass term. Taking  $H = 0$ , one obtains a linearized equation analogous to (4) that implies a growth rate  $\tau_{R,\text{eff}}^{-1} \propto m_\sigma^{\text{eff}} \propto (T_c^2 - T^2)^{1/2}$ . Growth is very slow near the  $T_c \approx \sqrt{2}v$ , prolonging the time over which the system is unstable. Only as the potential approaches its free space shape does the roll-down become rapid.

The annealing scenario describes the nonequilibrium evolution of the system in a regime where the time scale  $\tau_{\text{ev}}$  for the evolution of  $V_{\text{eff}}$  is much longer than  $\tau_R$ . One expects such a slow cooling scenario to further enhance the domain size — in metallurgy, the more rapidly one quenches, the smaller are the resulting crystals. The evolution of  $V_{\text{eff}}$  is determined by the expansion of the surrounding quasiparticles, which constitute a time-varying nonequilibrium heat bath. The annealing regime is relevant at RHIC energies because the temperature falls to  $T_c \approx 140$  MeV only at very late times, perhaps  $\tau_c \sim 10 - 20$  fm. The expansion time  $\tau_{\text{ev}}$  is then  $\sim \tau_c \gg \tau_R \sim 0.5$  fm.



**Figure 1:** The effective potential for various temperatures.

Müller and I explored the evolution in the annealing regime using a nonequilibrium Hartree-like approximation. [See Ref. [5] for details.] We indeed find that annealing prolongs the time during which the system is unstable compared to a quench. Domain sizes up to 7 fm are possible if we assume that  $m_\sigma^{\text{eff}}(T_c) = 0$  (*i.e.*, a second order transition). More conservatively, we find sizes of 3–4 fm for  $m_\sigma^{\text{eff}}(T_c) = 300$  MeV. Such large masses are possible, *e.g.*, in the large  $N$  limit of the  $O(N)$  model studied by Boyanovsky *et al.*<sup>10</sup> and Kluger *et al.*<sup>11</sup> However, QCD may be far richer than the large  $N$  limit implies.

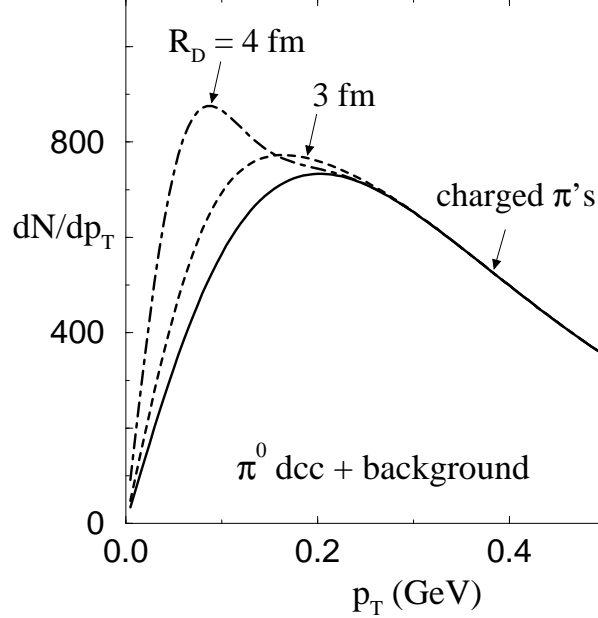
Suppose that a nuclear collision produces a large domain, either in a typical event due to the annealing dynamics or due to a rare fluctuation. How can we tell? The number of pions is not large, even ideally. The amount of energy available for pion production per unit volume is the potential difference  $\Delta V \sim 60 \text{ MeV fm}^{-3}$  between  $\Phi = 0$  and  $\Phi = (f_\pi, \vec{0})$ . If all of the energy in a domain of size  $R_D \sim 3\text{--}6$  fm goes into pion production, one expects  $N_{\text{dcc}} \sim \Delta V R_D^3 / m_\pi \sim 30\text{--}300$  pions. In comparison, conventional production mechanisms in a central Au+Au collision at RHIC can produce 1000 pions *per unit rapidity*; this constitutes a background to the DCC signal.

The source of pions from this domain is coherent and, therefore, concentrated at low transverse momenta. Horn and Silver, Gyulassy *et al.*, and Amado *et al.* have developed a coherent state formalism applicable to DCC production.<sup>13</sup> The amplitude for emitting a single neutral pion from a space time point  $(x_j, t_j)$  within a domain is real and satisfies  $A_j \propto \pi^0(x_j) \propto \exp\{-x_j^2/2R_D^2\}$ . The production rate from a single domain is

$$E \left( dN/d^3p \right)_{\text{dcc}} \propto \left| \sum_j A_j e^{-ip \cdot x_j + iE_{\vec{p}} t_j} \right|^2 \sim e^{-p^2 R_D^2}, \quad (7)$$

neglecting an energy-dependent factor that is slowly varying for small momenta. The rate is largest for  $p_T < R_D^{-1}$ .

The DCC signal (6) appears atop an isospin symmetric background. Assuming that all the DCC pions are neutral and taking the background produced by the lund/fritiof event generator for central Au+Au, I obtain the spectra in Fig. 2. A signal that is significant compared to statistical fluctuations in the background emerges at low  $p_T$  for a domain of size  $R_D > 3$ .



**Figure 2:** Neutral pion production from a single large DCC in Au+Au at RHIC.

Experimentally, one must search for DCCs on an event-by-event basis, looking for significant fluctuations in the neutral and charged pions. A domain producing a neutral fraction  $f$  introduces a contribution  $fE(dN/d^3p)_{\text{dcc}}$  to the  $\pi^0$  spectrum, with a corresponding contribution  $\{(1-f)/2\}E(dN/d^3p)_{\text{dcc}}$  to the  $\pi^+$  and  $\pi^-$  spectra. The distribution of  $f$  in a sample of large-domain events is given by (1).

To check that the pion emission from a candidate large-domain event is coherent, one can study the Bose correlations of pairs of identical pions. Pions produced by conventional scattering mechanisms are largely incoherent, and exhibit an intensity interference that is analogous to the Hanbury-Brown-Twiss effect for photons.<sup>15</sup> Measurements of  $\pi\pi$  correlations are typically used to study information on the size of the interaction volume in nuclear collisions, much as the HBT effect is used to measure the size of stars (although interpretation of experiments is highly nontrivial<sup>16</sup>). In contrast, a fully coherent source of pions exhibits interference at the amplitude level. Consequently, the single particle spectrum (7) from an isolated domain depends on the size of the domain, while the pair distribution is simply the product of single particle spectra.

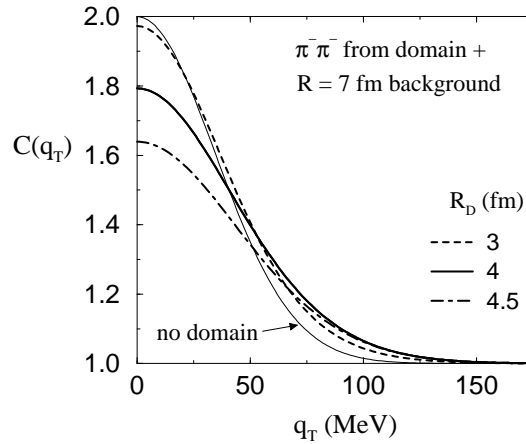
We therefore expect the HBT effect to be suppressed in a nuclear collision if a large domain forms. However, the HBT effect does not completely disappear, since intensity

interference can occur when a pion from the DCC has a similar momentum to a background pion. To estimate this effect, I follow Gyulassy *et al.* in Ref. [13] and write the two pion correlation function

$$C(p_1, p_2) = 1 + (1 - D(p_1))(1 - D(p_2))\tilde{\rho}(q)^2 + 2\{D(p_1)D(p_2)(1 - D(p_1))(1 - D(p_2))\}^{1/2}\tilde{\rho}(q), \quad (8)$$

where  $q = p_1 - p_2$  is the relative momentum of the pair and  $\tilde{\rho}$  is the Fourier transform of the space-time density of the background pions. Observe that (8) depends strongly on the fraction of coherent pions  $D = (dN/d^3p)_{\text{dcc}}/(dN/d^3p)_{\text{tot}}$ .

The possible impact of a DCC on  $\pi^-\pi^-$  correlations is illustrated in Fig. 3 as a function of domain size. I assume that the domain produces only charged pions and that the transverse source size for the fritiof background is 7 fm. A combination of one- and two-particle measurements can allow one to disentangle the DCC signal from the background. Of course, the interpretation of experiments will be extremely tricky!



**Figure 3:**  $\pi^-\pi^-$  correlations as a function of  $q_T = p_{T1} - p_{T2}$  from a single large DCC plus a background of incoherent pions in Au+Au at RHIC. The total momentum of the pair is taken to be  $K = (p_{T1} + p_{T2})/2 \equiv 0$ .

The schematic results in Figs. 2 and 3 indicate that domains larger than 3 fm can have measurable consequences. If seen in experiments, what would they tell us about QCD? All of these signals — the isospin fluctuations, the enhancement of the pion spectrum at low  $p_T$ , and the suppression of HBT correlations — are characteristics of any large coherent source. In principle, Bose condensation in an ideal pion gas can produce similar signals.<sup>16</sup> While it is argued<sup>17</sup> that the conditions at RHIC are not favorable for Bose condensation in the absence of a potential such as (3), systematic experimentation will nevertheless be needed to prove that the coherence comes from a disorientation of the chiral condensate. If an annealing scenario is valid, I expect large domains to occur in most central collisions.

Lego plots<sup>2</sup> could then reveal a ‘clumpy’ event structure indicative of different domains. The pion excess within each domain would be independent of the background multiplicity. On the other hand, the number of condensed pions in ideal Bose condensation would grow with the overall multiplicity and have a homogeneous structure.

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## References

1. K. Rajagopal and F. Wilczek, Nucl. Phys. **B379**, 395 (1993); **B404**, 577 (1993).
2. J.D. Bjorken, K.L. Kowalski and C.C. Taylor, SLAC preprint SLAC-PUB-6109 (April 1993); K.L. Kowalski and C.C. Taylor, Case Western Reserve Univ. preprint CWRUTH-92-6 (June, 1992), hep-ph/9211282; A.A. Anselm and M.G. Ryskin, Phys. Lett. **B266**, 482 (1989); J.-P. Blaizot and A. Krzywicki, Phys. Rev. **D46**, 246 (1992).
3. C. M. G. Lattes, Y. Fujimoto, and S. Hasegawa, Phys. Rep. **65**, 151 (1980); G. Arnison *et al.*, Phys. Lett. **122B**, 189 (1983); G. J. Alner *et al.*, Phys. Lett. **180B**, 415 (1986), Phys. Rep. **154**, 247 (1987); J. R. Ren *et al.*, Phys. Rev. **D 38**, 1417 (1988); Chacaltaya Collaboration, in proceedings of the *22nd International Cosmic Ray Conference* (Dublin, Ireland) **4**, 89 (1991); L. T. Baradzei *et al.*, Nucl. Phys. **B370**, 365 (1992).
4. S. Gavin, A. Gocksch and R.D. Pisarski, Phys. Rev. Lett. **72**, 2143 (1994).
5. S. Gavin and B. Müller, Phys. Lett. **B329**, 486 (1994).
6. S. Gavin, H. Satz, R.L. Thews, and R. Vogt, Z. Phys. **C61** (1994) 351; Nucl. Phys. **A566**, (1994) 383c; J.-P. Blaizot, and J.-Y. Ollitrault, in ‘Quark-Gluon Plasma,’ R. C. Hwa, ed. (World Scientific, Singapore, 1990), pp. 631-663.
7. F. R. Brown *et al.*, Phys. Rev. Lett. **65**, 2491 (1990); C. Bernard *et al.*, AZPH-TH-93-29, (1993); S. Gavin, A. Gocksch and R.D. Pisarski, Phys. Rev. **D49**, 307, (1994).
8. L. D. McLerran, Rev. Mod. Phys. **58** ( 1986) 1021-1064; J. Cleymans, R.V. Gavai, and E. Suhonen, Phys. Rept. **130**, 217 (1986).
9. J.F. Donoghue, E. Golowich, and B.R. Holstein, ‘Dynamics of the Standard Model,’ (Cambridge University Press, 1992).
10. D. Boyanovsky, H.J. de Vega, R. Holman, PITT-94-01 (1994), hep-ph/9401308.
11. F. Cooper, Y. Kluger, E. Mottola, J.P. Paz, Los Alamos Preprint (1994), hep-ph/9404357.
12. G. Baym and G. Grinstein, Phys. Rev. **D15** 2897 (1977); K. M. Benson, J. Bernstein, and S. Dodelson, Phys. Rev. **D44**, 2480, (1991).
13. D. Horn and R. Silver, Ann. Phys. (N.Y.) **66**, 509 (1971); M. Gyulassy, S.K. Kauffmann, L.W. Wilson, Phys. Rev. **C20** 2267 (1979); R.D. Amado, F. Cannata, J.P. Dedonder, M.P. Locher, Bin Shao, Phys. Rev. Lett. **72**, 970 (1994).
14. See, *e.g.* W. Zajc, in NATO Workshop on Particle Production in Highly Excited Matter, H. H. Gutbrod and J. Rafelski eds., (Plenum, New York, 1993), p. 427.



15. M. Gyulassy and S. S. Padula, Phys. Lett. **B217**, 181 (1989).
16. S. Pratt, Phys. Lett. **301**, 159, (1993).
17. C. Greiner, C. Gong and B. Müller, Phys. Lett. **B316**, 226 (1993).